Chapter 5

Diagnostic tools and techniques

The experimental observations and results presented in this thesis have been possible thanks to use of advanced and powerful diagnostic tools. The study of laser-plasma interactions presents a formidable challenge to present-day technology. In fact the phenomena studied in this thesis have typical scalelengths ranging from a fraction of a millimeter to a micron. Moreover the time resolution necessary to describe the evolution of the interaction was of the order of 1 to 100 picoseconds. For the experimental studies different diagnostic techniques have been used. This chapter presents a review of the main diagnostics used in the experiments.

5.1 Point-projection X-ray absorption spectroscopy

This powerful and versatile technique combines an imaging system of the target with information on the density and temperature of the plasma sample in the line of sight. Following earlier work in this field [Hoarty et al., 1997, 1999], the diagnostic was employed in the observation of laser-generated shock waves in low density foams, as reported in chapter 6.

A point-projection X-ray system consists of an X-ray point source which produces a flash of radiation. The X-rays are used to backlight the target (e.g. a plasma
Figure 5.1: Point-projection imaging system. The intensity profile on screen is drawn on the right-hand side.

sample) and are recorded on X-ray sensitive film. In our experiments the point source was generated by irradiating the tip of a gold pin with a short laser pulse. Figure 5.1 presents a schematic diagram of the point-projection imaging technique. The geometric magnification $M$ is given by

$$M = \frac{s + L}{s} = 1 + \frac{L}{s}, \quad (5.1)$$

where $s$ is the source-to-target distance, and $L$ is the target-to-film distance.

The spatial resolution can be estimated from obscuration tests. An object with a sharp edge and opaque to X-ray radiation is placed in the emission cone of the source, and its shadow is recorded on X-ray film. The intensity profile on screen would be a step profile in the case of an ideal point-source. The finite extent of the actual source produces a smoothed intensity profile, as shown in figure 5.1. The transition width $\delta x$ on screen relayed to the target plane can be used to estimate the resolution $\delta t$ of the imaging system (Rayleigh criterion)

$$\delta t = \frac{\delta x}{M}. \quad (5.2)$$

Here $\delta t$ is the minimum distance of two points in the target plane that result in separate images on the film plane. Given the relation $\frac{\delta x}{L} = \frac{\delta s}{s}$, where $\delta s$ is the source diameter, the resolution on target can be expressed as

$$\delta t = \delta s \frac{M - 1}{M} \sim \delta s \quad \text{for} \quad M \gg 1. \quad (5.3)$$
Hence it can be seen that for magnification factors $M$ greater than 10, the dimension of the source corresponds to the smallest resolved features on the target plane. Conversely, the transition width $\delta x$ can be used to estimate the effective source size, namely $\delta s = \delta x / (M - 1)$.

The actual set-up used in the work presented in the thesis combined the point-projection imaging system just described with an X-ray spectrometer. This addition allowed the separation on the film plane of the spectral components of the radiation coming from and through the plasma sample. Figure 5.3 shows a cross-sectional diagram of the spectrometer used in the experiments. The X-ray point source was placed in front of the spectrometer’s aperture. Some of the X-rays emitted from this backlighter probed the plasma and entered the spectrometer through a light-tight Be filter. After reflection from the crystal, the X-rays were recorded on film. An annular magnet deflected charged particles coming from the plasma to avoid damaging of the crystal inside the spectrometer.

By rotating the plane of the crystal it was possible to select the spectral range recorded on film.

A flat crystal separates spectral components according to Bragg’s law

$$n\lambda = 2d \sin \theta_B ,$$

where $n$ is the order of interference, $\lambda$ is the wavelength, $2d$ is twice the spacing between planes in the crystal, and $\theta_B$ is the Bragg’s angle. This angle corresponds to constructive interference for a particular wavelength. Each atom acts as a scattering centre for the incident beam (see figure 5.2). The difference in path length between rays emerging from planes in the crystal is a multiple of $2d \sin \theta_B$. Therefore constructive interference is obtained for the wavelength that satisfies Bragg’s law, whereas contribution from other wavelengths average to zero due to the large number of scattering planes.

The Rubidium Acid Phthalate (RbAP) crystal used in the experiment has a 2$d$ spacing of 26.12 Å, which places an upper limit upon the measurable wavelength of
λ = 26.12 Å. In the experiment the crystal was tilted to an angle of about 9° to the incoming radiation, in order to record on film the wavelength range from 3.2 to 5.6 Å (from 2.2 to 3.8 keV).

A RbAP crystal has a resolving power $E/\Delta E \sim 850$ [Burek, 1976], which corresponds to 3 eV resolution at 2.6 keV. This value must be corrected for the finite extension of the source.

Figure 5.2: Scattering from crystal planes.

Figure 5.3: Spectrometer (section).
5.2 Optical probing

The diagnostics based on optical probing comprise a vast number of different techniques, all of which use an independent laser beam to perform measurements of the plasma produced during the main laser interaction.

The presence of plasma in the line of propagation of this reference beam, also called \textit{probe beam}, is diagnosed by detecting the modifications (e.g. phase variations or angular deflections) induced on the beam. The term \textit{optical} probing, as opposed to X-ray probing, suggests that usually the probe wavelength is in or near the visible region of the spectrum. In general, higher harmonics of the interaction frequency are used to reduce the noise coming from the laser pulse.

In our experiments we used the 4th harmonic for better image contrast and higher depth of probing. The probe pulse was obtained by splitting a fraction of the uncompressed interaction pulse. As it will be seen in section 5.2.2, the probe and the interaction beam could be synchronised at the centre of interaction to within a few picoseconds, because both beams originated from the same pulse. After passing through the amplification stage, the probe beam was recompressed in a pair of gratings to a duration of 1-2 ps. The probe beam was converted into the fourth harmonic by passing through a type I KDP crystal and then through a type II KDP crystal, resulting in a wavelength of 266 nm.

The set-up used in the experiment is depicted in figure 5.4. The collimated probe beam entered and exited the interaction chamber through fused silica windows, which transmit UV radiation. After passing through the plasma, the probe was collected by an f/4 UV microscope objective which imaged the target plane on a screen outside the vacuum chamber. The best-focus position for the objective was checked by placing a 50 $\mu$m wire in the interaction centre and by taking a series of exposures corresponding to different objective-to-wire distances. The screen position was established in order to obtain a magnification on film of about 50 times.

An aperture was placed in the Fourier plane of the lens and acted as a frequency spatial filter, letting through only a narrow band of frequencies around the probe
5.2 Optical probing

Figure 5.4: Set-up for optical probing with a collimated probe beam.

frequency. Clearly the aperture was effective at filtering only rays almost parallel to the probe. In order to suppress the plasma self-emission radiation that went through the spatial filter, an interference filter, with passing bandwidth centred on the probe wavelength, was used to significantly reduce the intensity of the self-emission on film.

5.2.1 Interferometry

This technique is mainly used to measure density profiles and density modulations in a plasma. It is based on the phase map alterations induced by changes of the refractive index caused by the plasma. The relation that links the refractive index $\eta$ to the electron density and the critical density reads

$$\eta = \sqrt{1 - \frac{n_e}{n_c}}. \quad (5.5)$$

A probe beam going through a plasma acquires a phase difference $\Delta \phi$ with respect to a beam that has travelled the same distance $L$ in vacuum as a result of the modifications in the refractive index. The phase difference across the probe beam is
5.2 Optical probing

Figure 5.5: Cylindrical symmetric electron density profile and the plane section used in the evaluation of the phase shift integral.

given by the following integral along the propagation direction $y$ of the probe

$$
\Delta \phi(x, z) = -\frac{2\pi}{\lambda} \int_L dy \left( \eta(x, y, z) - 1 \right)
$$

Restricting our analysis to regions where the electron density is much smaller than the critical density $n_e \ll n_c$, the refractive index can be approximated by $\eta \approx 1 - \frac{n_e}{2n_c}$, and the phase integral is reduced to

$$
\Delta \phi(x, z) = \frac{\pi}{\lambda n_c} \int_L dy \, n_e(x, y, z)
$$

In cases where the electron density exhibits a cylindrical symmetry of the type $n_e(x, y, z) = n_e(r, z)$, it is possible to invert [Barr, 1962] the integral analytically, and $n_e$ can be expressed as a function of the phase shift. Referring to diagram 5.5 for the interpretation of the geometry, the phase shift integral 5.7 in cylindrical coordinates becomes

$$
\Delta \phi(x, z) = \frac{2\pi}{\lambda n_e} \int_{r=x}^{r=r_0} \frac{n_e(r, z)}{\sqrt{r^2 - x^2}} r \, dr,
$$

where $r_0$ represents the plasma boundary ($n_e \neq 0$) for each $z$ plane considered.

Equation 5.8 is one form of Abel’s integral equation, and it can be analytically
inverted under the assumption that $n_e(r, z) = 0$ for $r > r_0$ giving the electron density [Barr, 1962]

$$n_e(r, z) = -\frac{\lambda n_e}{\pi^2} \int_r^{r_0} \frac{\partial \Delta \phi(x, z)}{\partial x} \frac{dx}{\sqrt{x^2 - r^2}}.$$ (5.9)

Observing the symmetry of collected interferograms and comparing them with 2D simulations of the plasma expansion in our experimental conditions, we can state with confidence that all the assumptions made in deriving equation 5.9 are fulfilled in our experiments, and one is left with the task of calculating the phase integral with numerical methods.

In the data analysis we adopted the algorithm described in [Barr, 1962], which has the advantage of being computationally fast and relatively insensitive to small random errors in the data, since it has a built-in smoothing procedure. Once the phase distribution has been obtained from the interferogram in a matrix form $\Delta \phi(x_i, z_j)$ (see section 5.2.1), the axis of symmetry is identified and the phase is resampled in 21 points uniformly spaced in the outward radial direction starting from $r = 0$. Each row of the $\Delta \phi$ matrix is then multiplied by the 21x21 $\beta_{kn}$ matrix of the Barr coefficients, which performs the smoothing on the data and at the same time gives as a result the Abel inversion of the phase distribution.

**Set-up for interferometric measurements**

In the experiments reported in chapters 7 and 8 an interferometer was used to measure the density profile of the plasma at different times during the laser interaction. The interferometer was set up in a Nomarski modified configuration [Benattar et al., 1979, Borghesi et al., 1998], in which a birefringent Wollaston prism is used to generate interference between two parts of the probe beam. The experimental arrangement is depicted in figure 5.6. The phase object O (in our case the plasma) is relayed by the lens L on the film plane. The Wollaston prism W is set after the Fourier plane of L and splits the probe beam into two orthogonally polarised beams, that appear to be coming from two virtual foci F’ and F”.

Setting appropriately the
distance of the film plane from the prism, the two orthogonally polarised images can be made to overlap in a central region. The polariser $P_1$ selects a linearly polarised component of the probe. The polariser $P_2$ is rotated to 45 degrees with respect to each polarised beam that leaves the Wollaston prism, so that on the film plane two images of equal intensity and polarisation are produced. In the overlapping region a fringe pattern is generated with inter-fringe separation $\Delta x$ given by

$$\Delta x = \frac{\lambda p}{\epsilon q}.$$  \hfill (5.10)

where $\lambda$ is the probe wavelength, $p$ and $q$ are the distance of the Wollaston from the Fourier plane of L and from the image plane respectively. The aperture $\epsilon$ of the prism used in the set-up was about $10^{-3}$ rad. In the experiments the probe cross-section was larger than the plasma to be diagnosed, so that the image of the phase object could be made to interfere with an unperturbed part of the beam, as shown in figure 5.6. The fringe shift produced this way is exactly the one corresponding to equation 5.6.

It is essential for the good quality of the interferograms that the probe cross-section has a high degree of uniformity, since different parts of it are used to generate the interference pattern.
5.2 Optical probing

From the interferogram a) to the manually traced fringes b) and finally to the digitised and smoothed fringe pattern c). Lower panels present lineouts extracted from the corresponding images as indicated by the dashed line in panel a).

**Phase extraction**

The information about the phase modulations induced by the plasma is recorded on an interferogram as a pattern of fringes shifted from the original position. Together with it, there is the noise caused by the imaging optics and by the intrinsic spatial nonuniformity of the probe beam. Several computational techniques are reported in literature that are used to filter out the background noise and to retain only the phase information. In our analysis we found that still the most reliable and accurate approach was to trace manually the fringes on a transparent plastic foil, and digitise then the traced pattern with a high resolution scanner. The so-obtained fringes had a step function profile, which would cause the generation of spurious high frequencies in the subsequent analysis based on Fourier transform methods. Therefore the cos-squared profile of the actual fringes was recreated interpolating a sine function between each two adjacent fringes. Figure 5.7 shows this part of the data extraction.

Once the fringe pattern has been obtained, the phase shift can be determined using
5.2 Optical probing

a filter in the spatial frequency space as described in [Takeda et al., 1982, Gizzi et al., 1994]. Having oriented the unperturbed fringes along the x direction, the intensity \( I(x, z) \) of the fringe pattern can be written in the form

\[
I(x, z) = a(x, z) + b(x, z) \cos(2\pi f_0 z + \Delta \phi(x, z)) ,
\]

(5.11)

where \( a(x, z) \) and \( b(x, z) \) represent unwanted intensity variations across the probe beam, \( f_0 \) is the spatial-carrier frequency and \( \Delta \phi \) is the phase information we want to extract. It is important that the \( a \) and \( b \) nonuniformities have a spatial frequency lower than \( f_0 \), so that they can be decoupled from \( f_0 \) and filtered out in the frequency domain. Tracing by hand the fringes drastically reduces the contribution from such nonuniformities, effectively rejecting most of the noise and spurious signal.

Using the definition \( c(x, z) = \frac{1}{2} b(x, z) \exp(i\Delta \phi(x, z)) \), we can rewrite 5.11 as

\[
I(x, z) = a(x, z) + c(x, z)e^{i2\pi f_0 z} + c^*(x, z)e^{-i2\pi f_0 z} ,
\]

(5.12)

and, performing the Fourier transform along the \( z \) direction only, we obtain

\[
\hat{I}(x, f) = \hat{a}(x, f) + \hat{c}(x, f - f_0) + \hat{c}^*(x, f + f_0) .
\]

(5.13)

Here \( \hat{c}(x, f - f_0) \) represents the Fourier transform of \( c(x, z) \) translated by \( f_0 \). This part of equation 5.13 contains all the information we need. The method consists in selecting the part of the frequency spectrum that pertains to \( \hat{c}(x, f - f_0) \), shifting it back to the zero frequency by performing a translation by \( -f_0 \), and finally computing the inverse Fourier transform to obtain \( c(x, z) \). Algorithmically the whole process can be summarised as

\[
I(x, z) \mapsto \hat{I}(x, f) \mapsto \hat{c}(x, f - f_0) \mapsto \hat{c}(x, f) \mapsto c(x, z) .
\]

(5.14)

The last step consists in computing the logarithm of \( c(x, z) \) and retaining the imaginary part of it, which is the phase modulation we were looking for

\[
\ln[c(x, z)] = \ln[\frac{1}{2} b(x, z)] + i \Delta \phi(x, z) .
\]

(5.15)

As a final remark, we should point out that the so obtained phase \( \Delta \phi \) is in the range of \([0, 2\pi]\) as a consequence of the Fast Fourier Transform (FFT) algorithm used in the numerical analysis. Therefore a phase unwrapping procedure [Takeda et al., 1982] has to be implemented to recover the full ramp of the phase shifts.
Limit and sensitivity of the interferogram analysis

Refractive effects seriously limit the maximum electron density $n_e$ that can be probed by interferometry in our set-up. The upper limit is ideally set by the critical density $n_c$ for the probe wavelength, which in our case is of the order of $1.7 \cdot 10^{22} \text{ cm}^{-3}$. If the density gradient were only parallel to the probe, then densities close to $n_c$ could be measured. But when $\nabla n_e$ has a component transverse to the probe, as in any cylindrical symmetric plasma plume, the rays will be refracted out of the high density region. As soon as the refraction angle is greater than the acceptance of the imaging optics, the emerging rays are not collected and their information is lost. Ray-tracing simulations [Borghesi, 1998] for a probe propagating through a density profile similar to the ones produced in the experiments, show that for a probe wavelength $\lambda = 266 \text{ nm}$ and an objective with f-number of 4, rays are deflected out of the objective cone at densities of the order of $n_e = 1 - 2 \cdot 10^{20} \text{ cm}^{-3}$, i.e. about 2-5% of the critical density of the plasma forming beam. The previous considerations do not take into account any probe-plasma interaction such as absorption or scattering mechanisms. The probe beam will be affected in regions where the electron density is large enough for these processes to be significant. In our analysis we found that the maximum measured $n_e$ was of the order stated above for pure refraction effects. This indicates that the probe signal collected by the imaging optics was due to rays
refracted by the plasma.

The sensitivity of the interferometry technique is defined by the minimum detectable fringe shift $\delta \phi$. Using the Fourier transform method outlined in the previous section, the limiting factor that determines the minimum observable fringe shift is the small scale noise of the interferogram. A reliable estimate [Gizzi et al., 1994] for the minimum fringe shift in good quality interferograms is $\delta \phi = 2\pi/10 = 0.5$ radians, i.e. a tenth of the fringe separation. The corresponding minimum detectable electron density reads

$$\frac{\delta n_e}{n_e} = \frac{\delta \phi \lambda}{\pi L},$$

which in our experimental conditions gives $\delta n_e \sim 5 \times 10^{-4} n_e$ for a plasma scalelength $L$ of the order of 100 $\mu$m.

### 5.2.2 Timing

When collecting observations of a rapid varying phenomenon such as the laser-plasma interaction, it is essential to synchronise the diagnostic instruments with the interaction pulse on a timescale shorter than or at least comparable to the interaction timescale.

In part of the experiments presented in this thesis, the long-pulse beam lines of the VULCAN laser were used to generate the plasma, and then, at a suitable delay, the plasma was probed with a laser pulse and/or a proton beam. The laser probe and the particle beam were both generated using the short-pulse lines of the laser system. The two lines (short and long) are independent, and synchronisation between the two is achieved by using a Pockel cell gate system. The unavoidable jitter of this configuration is about 100 ps. To synchronise the interaction and the probe pulse, an optical streak camera was used with a resolution of a few tens of picoseconds. A piece of scattering material (in this case simply a piece of paper) was aligned in the centre of interaction of the vacuum chamber, and the scattered light was sent to the entrance slit of the streak camera. The beams were then fired in pairs and their optical path was changed using timing slides until the respective streaked images
overlapped.

**Figure 5.9:** Streak camera images illustrate the timing sequence of a 600 ps pulsed against a 1 ps pulse. Horizontal arrow indicates the direction of time. Small vertical arrow points to the 1 ps pulse signal. The short pulse comes a) early, b) after and c) on time with respect to the rising edge of the long pulse.

When a proton beam was used to probe a short-pulse interaction with a solid target, a synchronisation to within a few picoseconds became necessary. In this case the proton generating pulse, the probe pulse and the main interacting pulse were all split from the same pulse coming from the short-pulse line of VULCAN. Because they originated from the same seed pulse, the pulses could be timed to within their duration of about 1-2 ps by changing the optical path. A fast streak camera with 1 picosecond resolution was used to synchronise the three beams at the interaction centre.

### 5.3 Radiochromic film detector

In the following section the structure of the particle detector used in the experiments will be described along with the dosimetry properties and specifications of the detecting film.

The typical set-up for detecting and measuring a laser-generated proton beam is outlined in figure 5.10. The particle beam was produced by direct irradiation of a solid aluminium foil with an ultra-intense laser pulse. The beam is mainly composed of multi-MeV protons originated from a contamination layer of hydrocarbons.
The object to be probed was placed inside the propagation cone of the beam. The particles were then collected by the detector which had the capability of recording detailed information on the cross-section of the beam and on the energy distribution.

**Figure 5.10:** Typical set-up for proton probing used in the experiments. The proton beam emerged normal to the back surface of an aluminium foil after irradiation with an ultra-intense CPA laser pulse. The main target to be probed was placed in the propagation cone of the beam. Particles were collected in a detector consisting of a stack of radiochromic film.

The particle detector consisted of a stack of radiochromic film wrapped in aluminium foil. On the front surface, the Al foil acted as a soft X-ray filter and shielded the film from debris ejected from the main target after the laser interaction. High energy particles going through the film stack deposit locally their energy according to their stopping curve in the material. The radiochromic film changes colour upon exposure to ionising radiation, so each particle leaves a mark on the film along its track. In the geometry used in the experiments, the spatial distribution of the proton beam cross-section was recorded as an image on the radiochromic film. Each layer acted as an energy filter on the following layers, so that different energy components of the beam produced different signal on each film piece. Using this set-up, it was
possible to retrieve information on the energy distribution in the beam from the analysis and comparison of the layers.

### 5.3.1 Dosimetry specifications

The type of radiochromic film (RCF) used in the experiments presented in this thesis was the GAFCHROMIC® MD-55 produced by Nuclear Associates, USA. The active component is a monomer sensitive to ionising radiation. The monomer is in the form of sub-micron-sized crystals which are dispersed in a gelatin matrix and then coated onto a polyester film base [Associates, 2003]. The monomer molecule belongs to the family of diacetylenes, and when exposed to radiation in a crystalline state, it undergoes a solid-state polymerisation reaction producing a dye polymer referred to as polydiacetylene. The active crystals, which are transparent to visible light, change colour to a cyan blue when the dye is produced. The absorption spectrum of the dye is specific to the polydiacetylene molecule and for MD-55 dosimetry media has a major peak at 680 nm and a secondary peak at around 620 nm, as shown in figure 5.11. The amount of polymer produced, and therefore the depth of colour change, is proportional to the amount of energy absorbed in the active layer. After exposure, the film develops almost instantaneously and there is no need to treat the medium in any way to stop the reaction. The production of dye continues for a few hours while the reaction slows down, and the optical density stabilises to its asymptotic value within about 24 hours. The film is quite insensitive to room light, and does not need to be handled in the dark.

The response of the RCF MD-55 to ionising radiation is presented in figure 5.12. The change in net density\(^1\) is essentially a linear function of the dose, i.e. the energy deposited in the film, up to 50 Gy. The calibration data were obtained with

\(^1\)here *net density* is defined as the optical density of the exposed film minus the optical density of the unexposed film
5.3 Radiochromic film detector

Figure 5.11: Characteristic absorption spectrum of the dye polymer present in GAFCHROMIC® MD-55 medium, once exposed to radiation.

Figure 5.12: Response of GAFCHROMIC® MD-55 to Co$^{60}$ γ rays with visible broadband light source (dashed line) and with a 660 ± 10 nm filter (solid line).
γ rays from a Co$^{60}$ source. Dose fractioning effects are absent, as reported in the specification data sheet. The RCF is also dose rate independent in the range from 0.034 Gy/min to 3.4 Gy/min[Associates, 2003]. In our applications and analysis we made the assumption that the RCF had a dose independent response even at the extreme rates produced in the experiments, which were of the order of 10Gy/ps. Studies conducted on RCF irradiated with γ rays, electrons and protons show an identical response of the radiochromic medium to the various types of radiation [McLaughlin et al., 1991, Jones et al., 1999, Prasad and Bloch, 1995]. Regardless of the nature of ionising particles, the net density of the exposed film is a function of the amount of energy deposited in the active layers only. Saturation effects occur for the MD-55 film at doses of more than 100Gy, which were sometimes reached in the experiments here reported. The film was clearly over-radiated and the plastic base of the layers appeared to be visibly damaged.

The intrinsic resolution of the film is excellent, and allowed the observation of fine structures imprinted in the proton beam by the main target. In fact the polymerisation reaction initiated by radiation exposure does not spread between the microcrystals of which the active layer is composed, and therefore the film has intrinsically a sub-micron resolution.

![Figure 5.13: Structure and composition of a typical MD-55 radiochromic film.](image-url)
The structure of the MD-55 film is outlined in figure 5.13. An active layer, approximately 15 $\mu$m thick, is coated onto a 67 $\mu$m polyester (Mylar®) base. Two pieces of this film are then laminated coated-side to coated-side with a double-sided adhesive film. This adhesive film consists of a 25 $\mu$m thick polyester base coated on each side with a nominal 38 $\mu$m adhesive layer. A series of measurements gave for the thickness of the radiochromic film the average value of 268 $\pm$ 1 $\mu$m.

5.3.2 Energy deposition in the stack detector

When a high energy proton enters the RCF stack detector, it propagates through the film layers loosing energy primarily through ionisation of the medium molecules. Studies conducted on the form of the stopping curve of protons in RCF confirm that the response of the film is proportional to the Bragg curve of the protons in the radiochromic medium [Prasad and Bloch, 1995].

The Monte Carlo code SRIM [Ziegler et al., 2003] was used to simulate the energy deposition of a proton going through the detector. In the simulations, the detector was modelled by a 25 $\mu$m aluminium foil and a uniform block of polyester plastic with the same average density of the RCF. The chemical composition of the active layers is in fact quite similar to that of the plastic base. The code was run for a collimated monochromatic beam of protons hitting the detector at normal incidence. The number of protons traced through the detector was large enough to reach stabilisation of the stopping parameters, such as longitudinal range and lateral straggling, and of the ionisation curve as well. A series of stopping curves for different proton energies is shown in figure 5.14. The 25 $\mu$m Al foil in front of the RCF stack stops protons with energy less than 1.5 MeV. The presence of the Al filter is represented by the first peak at 25 $\mu$m of penetration. This peak is more evident at low energy, when the protons are almost stopped inside the filter.

Increasing the energy, the peak becomes more similar to a step. The Bragg peak depth into the detector is proportional to the proton energy. The value of the peak decreases as a consequence of the statistical spreading in energy distribution, which
**Figure 5.14**: SRIM simulation of the stopping curve of protons in the stack detector. The initial energy ranges from 2 MeV to 10 MeV at 1 MeV step. The stopping power is shown in keV/µm as a function of the penetration depth of the particles into the detector.

The stopping power becomes more important the longer the particle has travelled through the stopping medium. For this reason, the Bragg peak broadens significantly for higher energy particles. Several curves were simulated at 0.5 MeV steps in energy. Finally the curves were linearly interpolated to generate a *stopping surface*\(^2\) of the detector, as shown in figure 5.15. Using this surface, it was possible to estimate the stopping power for a proton at any point in the detector and for any initial energy in the range from 1.5 MeV to 20 MeV with continuity.

Given the typical structure of the RCF outlined in figure 5.13, the centres of the active layers are located in the detector at

\[
A_n = 25 \, \mu m + 268(n - 1) \, \mu m + 74.5 \, \mu m \quad \text{for } n = 1, 2, 3, \ldots
\]

and

\[
B_n = 25 \, \mu m + 268(n - 1) \, \mu m + 193.5 \, \mu m \quad \text{for } n = 1, 2, 3, \ldots
\]

\(^2\)By “stopping surface” we mean the surface generated by the stopping curves for a given pair \{particle,detector\} when the energy of the incident particle is varied continuously.
Figure 5.15: Simulated stopping surface for a RCF film stack detector. Values of the stopping power are given in keV/µm at any depth in the detector in µm, and for a range of initial proton energies expressed in MeV.
where \( n \) is the progressive number of RCF film piece, and the first and second active layers of each RCF film are labelled \( A \) and \( B \) respectively. Active layer \( A_1 \) is the closest to the aluminium foil filter.

For any given active layer, the energy deposition as a function of initial proton energy was obtained by integrating the stopping surface over the region spanned by the active layer. Figures 5.16 and 5.17 show the energy deposition in the first active layer \( A_1 \) and in the first RCF film respectively. The latter is the sum of the contributions from both active layers. This operation is consistent with the linear response of the RCF optical density to the dose. Therefore the optical density of the sandwiched radiochromic film is proportional to the sum of the doses absorbed in each active layer.

The energy distribution of protons is assumed to be uniform in these estimates, so that the curves indicate the energy deposited on average in the active layers for any given proton. When investigating the temperature of the proton beam, the contribution of each energy component was rescaled according to the assumed energy distribution, and the estimated response for each active layer was obtained. In our analysis the beam energy distribution was generally described by an exponential \( \exp(-E/T) \), where \( T \) represented the temperature parameter of the beam.
Figure 5.16: Energy deposition in the first active layer of the detector as a function of the initial proton energy.

Figure 5.17: Energy deposition in the first piece of radiochromic film, obtained by summing the response of the first and second active layer.
5.3 Radiochromic film detector

5.3.3 From proton energy to a time sequence

Since the stack detector operates a selection on the energy range recorded in each layer, it is possible to reconstruct a multi-frame sequence of a phenomenon evolving in time. For a given proton energy $E$, the probing time $t_{prob}$ for the particle can be estimated by the time of flight $t_{tof}(E)$ of the particle from the source to the target plane. The time of flight is readily calculated from the set-up geometry, and it reads

$$t_{tof}(E) = \frac{s}{v_p} = \frac{s}{c} \sqrt{\frac{m_p}{2E}}$$

(5.17)

where $s$ is the source-to-target distance, and $m_p$ and $E$ are both expressed in MeV. Protons in the ion beam will reach the target at different times, according to their energy. In the experiment presented in chapter 8, the ion beam was timed against a picosecond laser-solid interaction, and it was possible to observe the evolution of e.m. fields around the interaction region on a timescale of a few tens of picosecond. In that experiment the proton beam had an estimated temperature of 1.5 MeV and the resulting response of the stack detector is shown in figure 5.18. Using the

Figure 5.18: Fractional dose as a function of proton beam energy component. The temperature of the beam is set to $T = 1.5$ MeV. Curves for the first six active layers are shown. Each curve has been renormalised to its peak value. Labels on the top mark the active layers.
Figure 5.19: Signal recorded on each active layer as a function of the probing time of protons. Labels on the top mark the active layers.

Formula in 5.17 and the detector response to each energy component, it is possible to obtain the signal on each separate active layer as a function of the probing time. Figure 5.19 presents the results of this calculation for a source-to-target distance $s = 2$ mm. In this case, the maximum observable window amounts to about 30 ps, and the frame-to-frame separation ranges from 13 to 2 ps.